

Supersymmetric effects on the Forward Backward asymmetries of $B \rightarrow K\tau^+\tau^-$

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Leptonic and semi-leptonic rare decays of B-mesons provide significant (both theoretically and experimentally) signatures of any new physics beyond the Standard Model (SM). More specifically the decay $B \rightarrow K\ell^-\ell^+$ has been theoretically observed to be very sensitive to new physics, as the Forward Backward (FB) asymmetry in this decay mode vanishes in the SM. Supersymmetry, however, predicts a non-vanishing value of this asymmetry. In this work we will study the polarized lepton pair FB asymmetry, i.e. the FB asymmetry of the lepton when one (or both) final state lepton(s) are polarized. We will study these asymmetries both within the SM and for Supersymmetric corrections to the SM.

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I. INTRODUCTION

Lately there has been enormous progress in the study of flavour physics, where the B system has provided us with one of the most ideal environments for this type of study. Of the decay modes considered, theoretically and experimentally, the study of the “*rare decays*” of the B-mesons are of particular interest. Here the name *rare* has been given to those decay modes which arise from Flavour Changing Neutral Currents (FCNC). FCNC processes are absent at the tree level in the Standard Model (SM) but can occur through loop diagrams, their strength being proportional to the Fermi Constant. FCNC processes involving, for example $b \rightarrow s$ and $b \rightarrow d$ transitions are therefore more sensitive to the details of the SM interactions and thus are suited to the study of possible new physics beyond the SM. Theoretically inclusive FCNC processes like $B \rightarrow X_{s(d)}\ell^+\ell^-$ are relatively cleaner than their exclusive counterparts, since they are relatively independent of the quark structure of the hadrons involved. They are however difficult to measure experimentally (for details refer to Chapter 7 in reference [24]) and one may expect a substantial amount of experimental information regarding various exclusive B decay processes from the B-factories. Amongst the important FCNC exclusive B-decay processes are $B \rightarrow \text{charmless meson} + \gamma$ and $B \rightarrow \text{charmless meson} + \text{lepton pair}$. This second process potentially provides a very rich set of experimental observables involving various momenta and spin polarization correlations. It is thus important to theoretically calculate all possible measurable parameters of these processes.

In recent times there have been many calculations of such processes like $B \rightarrow K(K^*)\ell^+\ell^-$ [8, 9, 10], $B \rightarrow \pi(\rho)\ell^+\ell^-$ [13], $B_{s,d} \rightarrow \ell^+\ell^-$ [14] and $B_{s,d} \rightarrow \ell^+\ell^-\gamma$ [16]. Amongst these the ones involving a quark level $b \rightarrow s$ transition are expected to have relatively large branching ratios. For the $B \rightarrow K^*\ell^+\ell^-$ transition possible experimentally accessible parameters like Forward-Backward asymmetry, lepton polarization asymmetry etc., have been studied [2, 7, 8, 10]. In particular, polarization correlations between the two leptons, which was suggested recently by Bensalem *et al.*[5] have

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also been studied in the context of this process [6]. In this note we carry out an analysis of the polarized Forward Backward (FB) asymmetries in the process $B \rightarrow K\ell^+\ell^-$.

The theoretical basis for the study of FCNC B-decay processes is now a well established formalism based on the operator product expansion and use of the renormalization group [1]. The formalism ultimately produces an effective Hamiltonian for every process involving low dimensional hadronic operators, in the form of currents, with numerical multiplying coefficients called the Wilson coefficients. The Wilson coefficients encrypt short distance properties of the weak Hamiltonian and are sensitive to physics beyond the SM at high energy scales, in particular, to supersymmetric extensions of the SM. A great deal of theoretical work has gone, in recent times, to evaluating their values both within the SM and in the context of the minimal extension of the standard model (MSSM). Evaluation of the matrix elements of the hadronic currents on the other hand involve the relatively long distance quark structure of the hadrons. The most ideal way to evaluate them would be through lattice gauge theory calculations, which do not exist for all B-meson processes at the moment. Alternatively, one relies on evaluations based on Light cone sum rules and these also have been compiled for a variety of processes. We shall make extensive use of these in our calculations.

This paper is organized as follows: In section II we will discuss the effective Hamiltonian for the process under consideration. In section III we will introduce our notation for the polarized FB asymmetry. Finally, in section IV we will present results of our numerical analysis and the conclusions.

II. EFFECTIVE HAMILTONIAN

In this paper we are interested in the process $B \rightarrow K\ell^-\ell^+$, which has the basic quark level transition $b \rightarrow s$. The effective Hamiltonian for a such transition has been summarized in the literature [1]. The effective Hamiltonian is arrived at by integrating out the heavy degrees of freedom from the full theory. In the SM the heavy degrees of freedom are W^\pm, Z and the top quark; in the MSSM all the new SUSY particles shall also be counted. From such considerations we arrive at the effective Hamiltonian [17] :

$$\mathcal{H}_{eff} = \frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \left[\sum_{i=1}^{10} C_i(\mu) O_i(\mu) + \sum_{i=1}^{10} C_{Q_i}(\mu) Q_i(\mu) \right] \quad (2.1)$$

where the O_i are the current-current ($i = 1, 2$), penguin ($i = 3, \dots, 6$), magnetic penguin ($i = 7, 8$) and semi-leptonic ($i = 9, 10$) operators, whereas the $C_i(\mu)$ are the corresponding Wilson coefficients renormalized at scale μ . The value of these coefficients have been given in references [18, 19]. The additional operators Q_i ($i = 1, \dots, 10$), and their corresponding Wilson coefficients are due to the Neutral Higgs boson (NHB) exchange diagrams and are given in references [14, 17].

Different FCNC decays involve different combinations of C_i 's and C_{Q_i} 's and thus provide us with independent information on these coefficients. At the quark level the transition $b \rightarrow s\gamma$ is sensitive only to the magnitude of C_7 , whereas the semi-leptonic transition $b \rightarrow s\ell^-\ell^+$ is sensitive to C_9, C_{10}, C_{Q_1} and C_{Q_2} as well¹.

From the effective Hamiltonian given in equation (2.1) the decay amplitude for $B \rightarrow K\ell^-\ell^+$ is calculated to be:

$$\mathcal{M} = \frac{\alpha G_F}{\sqrt{2}\pi} V_{tb} V_{ts}^* \left\{ -2C_7^{eff} \frac{m_b}{q^2} (\bar{s}i\sigma_{\mu\nu} q^\nu P_R b)(\bar{\ell}\gamma^\mu \ell) + C_9^{eff} (\bar{s}\gamma_\mu P_L b)(\bar{\ell}\gamma^\mu \ell) + C_{10} (\bar{s}\gamma_\mu P_L b)(\bar{\ell}\gamma^\mu \gamma_5 \ell) + C_{Q_1} (\bar{s}P_R b)(\bar{\ell}\ell) + C_{Q_2} (\bar{s}P_R b)(\bar{\ell}\gamma_5 \ell) \right\} \quad (2.2)$$

where q is the momentum transferred to the lepton pair, given as $q = p_- + p_+$, where p_- and p_+ are the momentas of ℓ^- and ℓ^+ respectively. The $V_{tb}V_{ts}^*$ are the CKM factors and $P_{L,R} = (1 \mp \gamma_5)/2$. In writing the above matrix element (and in future analysis) we will neglect the mass of the strange quark, whereas lepton masses shall be retained.

The free quark decay amplitude given in equation (2.2) contains certain long distance effects which are absorbed in the redefinition of the C_9 Wilson coefficient (where we use the prescription given in reference [20]) :

$$C_9^{eff}(\hat{s}) = C_9 + Y(\hat{s}). \quad (2.3)$$

The $Y(\hat{s})$ part has a perturbative as well as a non-perturbative part. The origin of the non-perturbative part is from the resonance corrections to the perturbative quark loops (which gives the perturbative contribution to $Y(\hat{s})$). We

¹ It has also been shown in many works that $b \rightarrow s\ell^-\ell^+$ is sensitive even to the signs of these Wilson coefficients and hence this decay channel will provide information not only on the magnitude but also the sign of these coefficients.

will also use the usual Breit-Wigner prescription to take care of the resonant contribution [7, 20]. This prescription implies adding resonant terms to C_9^{eff} :

$$C_9^{res} \propto \kappa \sum_{V=\psi} \frac{\hat{m}_V Br(V \rightarrow \ell^-\ell^+) \hat{\Gamma}_{total}^V}{\hat{s} - \hat{m}_V^2 + i\hat{m}_V \hat{\Gamma}_{total}^V} \quad (2.4)$$

where all the symbols above have been explained in the work of Krüger and Sehgal [7]. For the phenomenological factor, κ , we will choose a value 2.3.

Using the definitions of the form factors given in Appendix A we can write the matrix element given in equation (2.2) as:

$$\mathcal{M} = \frac{\alpha G_F}{2\sqrt{2}\pi} V_{tb} V_{ts}^* \left\{ A(p_K)_\mu (\bar{\ell}\gamma^\mu \ell) + B(p_K)_\mu (\bar{\ell}\gamma^\mu \gamma_5 \ell) + C(\bar{\ell}\ell) + D(\bar{\ell}\gamma_5 \ell) \right\} \quad (2.5)$$

where the coefficients A, B, C and D in equation (2.11) are given as²:

$$A = 4f_T \frac{\hat{m}_b C_7^{eff}}{1 + \hat{m}_K} + 2f_+ C_9^{eff} \quad (2.6)$$

$$B = 2f_+ C_{10} \quad (2.7)$$

$$C = f_0 \frac{1 - \hat{m}_K^2}{\hat{m}_b} C_{Q_1} \quad (2.8)$$

$$D = f_0 \frac{1 - \hat{m}_K^2}{\hat{m}_b} C_{Q_2} + 2\hat{m}_\ell C_{10} f_+ - 2\hat{m}_\ell f_+ \frac{1 - \hat{m}_K^2}{\hat{s}} C_{10} + 2\hat{m}_\ell f_0 \frac{1 - \hat{m}_K^2}{\hat{s}} C_{10} \quad (2.9)$$

With the above expression of the matrix element, equation (2.5), we can obtain the expression for the differential decay rate as:

$$\frac{d\Gamma}{d\hat{s}} = \frac{\alpha^2 G_F^2 m_B^5}{2^{12} \pi^5} |V_{tb} V_{ts}^*|^2 \lambda^{1/2} \sqrt{1 - \frac{4\hat{m}_\ell^2}{\hat{s}}} \Delta \quad (2.10)$$

where

$$\begin{aligned} \Delta = & \frac{2}{3} \lambda \left(1 + \frac{2\hat{m}_\ell^2}{\hat{s}} \right) |A|^2 + \frac{2}{3} \left[\lambda \left(1 + \frac{2\hat{m}_\ell^2}{\hat{s}} \right) + 24\hat{m}_K^2 \hat{m}_\ell^2 \right] |B|^2 + 4(\hat{s} - 4\hat{m}_\ell^2) |C|^2 + 4\hat{s} |D|^2 \\ & + 8\hat{m}_\ell (1 - \hat{m}_K^2 - \hat{s}) \text{Re}(B^* D) \end{aligned} \quad (2.11)$$

with $\lambda = 1 + \hat{m}_K^4 + \hat{s}^2 - 2\hat{m}_K^2 \hat{s} - 2\hat{m}_K^2 - 2\hat{s}$. In the next section we will use this expression for the differential decay rate to introduce and define the polarized FB asymmetries, followed by our analytical expressions for these asymmetries.

III. POLARIZED FB ASYMMETRIES

Firstly we will define the polarization vectors of ℓ^- and ℓ^+ , where in this definition we will use the convention followed in many earlier works [6, 7, 8, 9]. In order to evaluate the polarized FB asymmetries we introduce a spin projection operator defined by $N = (1 + \gamma_5 S_x)/2$ for ℓ^- and $M = (1 + \gamma_5 W_x)/2$ for ℓ^+ , where $x = L, N$, or T (corresponding to the longitudinal, normal and transverse polarization asymmetries respectively). Next we define the orthogonal unit vectors S_x for ℓ^- and W_x for ℓ^+ in the rest frames of ℓ^- and ℓ^+ respectively as:

$$\begin{aligned} S_L^\mu &\equiv (0, \mathbf{e}_L) = \left(0, \frac{\mathbf{p}_-}{|\mathbf{p}_-|} \right) \\ S_N^\mu &\equiv (0, \mathbf{e}_N) = \left(0, \frac{\mathbf{p}_K \times \mathbf{p}_-}{|\mathbf{p}_K \times \mathbf{p}_-|} \right) \end{aligned}$$

² In writing equation (2.11) we have used the equations of motion: $p_\mu \bar{\ell} \gamma^\mu \ell = 0$; $p_\mu \bar{\ell} \gamma^\mu \gamma_5 \ell = 2m_\ell (\bar{\ell} \ell)$.

$$S_T^\mu \equiv (0, \mathbf{e}_T) = (0, \mathbf{e}_N \times \mathbf{e}_L) \quad (3.1)$$

$$\begin{aligned} W_L^\mu &\equiv (0, \mathbf{w}_L) = \left(0, \frac{\mathbf{p}_+}{|\mathbf{p}_+|}\right) \\ W_N^\mu &\equiv (0, \mathbf{w}_N) = \left(0, \frac{\mathbf{p}_K \times \mathbf{p}_+}{|\mathbf{p}_K \times \mathbf{p}_+|}\right) \\ W_T^\mu &\equiv (0, \mathbf{w}_T) = (0, \mathbf{w}_N \times \mathbf{w}_L) \end{aligned} \quad (3.2)$$

where \mathbf{p}_- , \mathbf{p}_+ and \mathbf{p}_K are the three momenta of ℓ^-, ℓ^+ and the K-meson in the dilepton CM frame. From the rest frames of the leptons we boost the four vectors S_x and W_x to the dilepton CM frame. Only the longitudinal vectors, S_L and W_L will be boosted by the Lorentz transformation, to a value of :

$$\begin{aligned} S_L^\mu &= \left(\frac{|\mathbf{p}_-|}{m_\ell}, \frac{E_\ell \mathbf{p}_-}{m_\ell |\mathbf{p}_-|}\right) \\ W_L^\mu &= \left(\frac{|\mathbf{p}_-|}{m_\ell}, -\frac{E_\ell \mathbf{p}_-}{m_\ell |\mathbf{p}_-|}\right) \end{aligned} \quad (3.3)$$

where E_ℓ is the energy of either of the leptons (both having the same energy in this frame) in the dileptonic CM frame.

The definition of the differential Forward-Backward (FB) asymmetry is given in references [2, 21]:

$$\overline{A}(\hat{s}) = \int_0^1 \frac{d^2\Gamma}{d\hat{s} dz} dz - \int_{-1}^0 \frac{d^2\Gamma}{d\hat{s} dz} dz. \quad (3.4)$$

Consider the case where we shall not sum over the spins of the outgoing leptons. In general the FB asymmetry will be a function of the spins of the final state leptons, and as such can be defined as :

$$\overline{A}(s^-, s^+, \hat{s}) = \int_0^1 \frac{d^2\Gamma(s^-, s^+)}{d\hat{s} dz} dz - \int_{-1}^0 \frac{d^2\Gamma(s^-, s^+)}{d\hat{s} dz} dz. \quad (3.5)$$

From an experimental viewpoint the normalized FB asymmetry is more useful. Therefore we shall normalize the above expression (equation 3.5) for the FB asymmetry by dividing by the total decay rate:

$$A(s^-, s^+, \hat{s}) = \frac{\overline{A}(s^-, s^+, \hat{s})}{d\Gamma/d\hat{s}}. \quad (3.6)$$

In analogy to the prescription given in Bensalem *et al.*[5] we can split this FB asymmetry into various polarization components³:

$$A(s^-, s^+) = A + A_i^- s_i^- + A_i^+ s_i^+ + A_{ij} s_i^+ s_j^- \quad (3.7)$$

where $i, j = L, T, N$ are the longitudinal, transverse and normal components of the polarization. Using this definition we can write the single and double lepton polarized FB asymmetries. From equation (3.7) the single polarized lepton FB asymmetry can be written as:

$$A_i^- = A(s^- = i, s^+ = j) + A(s^- = i, s^+ = -j) - A(s^- = -i, s^+ = j) - A(s^- = -i, s^+ = -j) \quad (3.8)$$

$$A_i^+ = A(s^- = j, s^+ = i) + A(s^- = -j, s^+ = i) - A(s^- = j, s^+ = -i) - A(s^- = -j, s^+ = -i) \quad (3.9)$$

and the doubly polarized FB asymmetry can be written as:

$$A_{ij} = A(s^- = i, s^+ = j) - A(s^- = i, s^+ = -j) - A(s^- = -i, s^+ = j) + A(s^- = -i, s^+ = -j). \quad (3.10)$$

Using the above expressions of the FB asymmetries the results of the unpolarized FB asymmetry is evaluated to be :

$$A = 2\hat{m}_\ell \lambda \sqrt{1 - \frac{4\hat{m}_\ell^2}{\hat{s}}} \frac{Re(A^*C)}{\Delta}. \quad (3.11)$$

³ The convention followed is that the repeated index is summed over.

From the expression given above and equation (2.8), we can see that the unpolarized FB asymmetry is proportional to C_{Q_1} . This point has been emphasized in many earlier works [12, 13]. In the SM C_{Q_1} is absent and hence for the decay modes $B \rightarrow K(\pi)\ell^+\ell^-$ the FB asymmetry within the SM vanishes. However, in SUSY (and 2HDM) extensions of the SM there exists a non-vanishing value of C_{Q_1} , and hence a non-vanishing value of the FB asymmetry [12, 13] is possible. Therefore a non-vanishing value of the FB asymmetry can be regarded as clear signal of new physics beyond the SM.

The analytical results of the polarized FB asymmetries are :

$$A_L^- = \frac{4\lambda^{1/2}\hat{m}_\ell}{\Delta} \left[\hat{m}_\ell(-1 + \hat{m}_K^2 + \hat{s}) \text{Re}(A^*B) - \hat{m}_\ell \text{Re}(A^*D) \right] \quad (3.12)$$

$$A_N^- = 0 \quad (3.13)$$

$$A_T^- = \frac{4\hat{m}_\ell}{3\hat{s}\Delta} \sqrt{1 - \frac{4\hat{m}_\ell^2}{\hat{s}}} \lambda \text{Re}(A^*B) \quad (3.14)$$

$$A_L^+ = A_L^- \quad (3.15)$$

$$A_N^+ = 0 \quad (3.16)$$

$$A_T^+ = A_T^- \quad (3.17)$$

$$A_{LL} = A \quad (3.18)$$

$$A_{LN} = \frac{4\hat{m}_\ell}{3\sqrt{\hat{s}}\Delta} \sqrt{1 - \frac{4\hat{m}_\ell^2}{\hat{s}}} \lambda \text{Im}(A^*B) \quad (3.19)$$

$$A_{LT} = \frac{4\hat{m}_\ell}{3\sqrt{\hat{s}}} \sqrt{1 - \frac{4\hat{m}_\ell^2}{\hat{s}}} \lambda \frac{|A|^2}{\Delta} \quad (3.20)$$

$$A_{NL} = A_{LN} \quad (3.21)$$

$$A_{NN} = -A \quad (3.22)$$

$$A_{NT} = 4\hat{m}_\ell \lambda^{1/2} \frac{(1 - \hat{m}_K^2 - \hat{s})\hat{m}_\ell \text{Im}(A^*B) + \text{Im}(A^*D)}{\Delta} \quad (3.23)$$

$$A_{TL} = -A_{LT} \quad (3.24)$$

$$A_{TN} = A_{NT} \quad (3.25)$$

$$A_{TT} = A \quad (3.26)$$

where Δ is given in equation (2.11) and A is the unpolarized FB asymmetry given in equation (3.11). We will discuss the above obtained expressions of the various FB asymmetries and present our numerical analysis of the same in the next section.

IV. NUMERICAL ANALYSIS, RESULTS AND CONCLUSION

In this section we shall present our numerical analysis of the observables whose analytical expressions were given in the previous section. We will also present the variation of all the observables with the dilepton invariant mass.

As it is experimentally more useful to have the average values of these quantities we shall present our results as the averages values of these quantities, where we will define our averages as:

$$\langle A \rangle \equiv \frac{\int_{(3.646+0.02)^2/m_B^2}^{(m_B-m_K)^2/m_B^2} A \frac{d\Gamma}{ds} d\hat{s}}{\int_{(3.646+0.02)^2/m_B^2}^{(m_B-m_K)^2/m_B^2} \frac{d\Gamma}{ds} d\hat{s}} \quad (4.1)$$

which means that in calculating the average we have taken the lower limit of integration to be above the first resonance⁴. The input parameters of our numerical analysis are defined in Appendix B, and our SM predictions of the integrated observables are given in Table I.

⁴ The first resonance here means the resonance after the threshold of the decay, which is $s \geq 4m_\tau^2$.

$\text{Br}(B \rightarrow K\tau^-\tau^+)$	A	A_L^-	A_T^-	A_{LN}	A_{LT}	A_{NT}
1.17×10^{-7}	0	0.363	-0.097	0.023	0.187	.0847

TABLE I: The SM predictions for the integrated observables

Before discussing our results we shall first elaborate on the models in which we have performed our numerical analysis. We have worked with the minimal supersymmetric extension of the standard model (MSSM), this being the simplest SUSY extension of the SM with the least number of parameters introduced [22]. But even in the MSSM we are required to introduce a large number of parameters, over and above the number of parameters in the SM. To ease out this problem and to reduce such a large number of parameters to a more manageable level, many models have been introduced such as the dilaton, moduli, minimal Supergravity (mSUGRA) [27], AMSB (Anomaly mediated SUSY breaking) [28] and the GMSB (Gauge mediated SUSY breaking) [29] models. The generic feature of all these models is that they assume some sort of unification of the parameters of the MSSM at a higher scale. In the literature the mSUGRA model is also known as the CMSSM (Constraint MSSM) [30]. We shall further assume that the soft SUSY breaking parameters are real. For our numerical analysis we will use two types of SUGRA (Supergravity) models, namely mSUGRA and rSUGRA (relaxed SUGRA) which we describe. In both these models it is believed that the SUSY breaking occurs in a hidden sector and is communicated to the visible sector only by gravitational-strength interactions. As such, soft breaking terms are assumed to be flavour blind (like gravitational interactions).

In the mSUGRA model along with the unification of the coupling constants $g_{1,2,3}$ (of the U(1), SU(2) and SU(3) gauge theories), the other unification conditions are;

- unification of the gaugino masses ($m_{1/2}$),
- unification of the scalar (sfermion and Higgs) masses (m_0),
- unification of the trilinear couplings (A),

all at the GUT scale. There are two other parameters. The first is the ratio of vev (vacuum expectation value of two Higgs), namely $\tan\beta$. The second arises in the process of evolving the soft SUSY breaking parameters from the GUT scale to the electroweak scale and then imposing the correct low energy electroweak symmetry breaking condition. This condition fixes the magnitude of μ (the two Higgs coupling parameter), however, the sign still remains uncertain⁵. The sign of μ thus also enters taken as a parameter.

Therefore, in all mSUGRA frameworks we have five parameters, namely:

$$m_{1/2}, m_0, A, \tan\beta \text{ and } \text{sgn}(\mu) \quad (4.2)$$

Unification of all the scalars and also all the gauginos is not an essential requirement of SUGRA models. One can have models where either all the scalars do not have a universal mass at GUT scale or there is a non-universality of gaugino masses at the GUT scale. We shall also explore such a model, where we would relax the condition of the universality of the scalar masses at the GUT scale [23]. In the literature these models are known as non minimal SUGRA models [23]. We will call such a model to be relaxed SUGRA (rSUGRA) model. We will assume that the values of squarks and the Higgs sector scalars masses different at the GUT scale. This shall introduce another parameter into the model. This additional parameter we will take as the mass of the pseudo-scalar Higgs boson (m_A). However for our numerical analysis we will consider only the region of the SUGRA parameter space which is consistent with the $B \rightarrow X_s\gamma$ 95% CL [3, 24]:

$$2 \times 10^{-4} \leq \text{Br}(B \rightarrow X_s\gamma) \leq 4.5 \times 10^{-4}.$$

We present our results for the various decay rates and asymmetry parameters considered in Figures (1)-(21). The branching fractions of the decays considered are of course too low to be observed with the current luminosities of the B-factories, but they will certainly be possible in the foreseeable future. As can be seen from the expressions of the polarized FB asymmetry, they are sensitive to the Wilson coefficients that arise only beyond the standard model, and thus a measurement of these would be one more test of physics beyond the SM.

⁵ There are many conventions followed regarding the sign of μ , our convention is such that μ appears in the chargino mass matrix with a positive sign (where the $(g - 2)_\mu$ giving the parameter μ a negative sign is disfavored).

We now turn to the uncertainties in the estimates we obtained, where the Wilson coefficients we used are in the NLL approximation and do not introduce significant uncertainties. The CKM parameters typically have uncertainties of the order of 10% [31] while the other SM parameters involved do not suffer from any large uncertainties. The SUSY parameters are input parameters. For a particular choice of parameters the major uncertainty in our results arises from the definition of the form factors used. They typically have a 15% uncertainty [26]. There are regions in the SUSY parameter space where SUSY effects exceed the kind of errors which are introduced by the form factor definition, for example the graph of the branching ratio of the mode concerned $B \rightarrow K\ell^-\ell^+$ in figures 1, 2, 3. Of course, at the present level an uncertainty of 20 % modifications arising from SUSY or any other physics beyond the SM, or levels less than this, would not be distinguishable from SM results. We, however, emphasize that the null polarization results which we have pointed out for some of the double polarization asymmetries (for a mixture of B_0 and \bar{B}_0) do not depend on the parameters. The experimental deviation of these from null values would indicate a presence of interactions beyond the one considered here, or may indicate the source of CP-violation in the SUSY extension, for example a complex value of μ .

Regarding our numerical analysis we have presented all the observables in section III, where the plots are given for all the possible observables with the dilepton invariant mass (\hat{s}), in figures 1, 4, 7, 10, 13, 16, 19 for the SM, mSUGRA and rSUGRA models. We have also tested the sensitivity of the observables to various MSSM parameters. For this purpose we have also presented the results of the averaged values of the observables. For averaging we have used the procedure defined in equation (4.1). The plots of averaged observables with $\tan\beta$ for various values of m_0 are given in figures 2, 5, 8, 11, 14, 17, 20 in the mSUGRA model. The other model parameters are given in respective figure captions. In the rSUGRA model we have plotted the averaged value of the observables as a function of pseudo-scalar Higgs mass (m_A) for various values of $\tan\beta$ in figures 3, 6, 9, 12, 15, 18, 21.

As can be seen from the plots all the observables are very sensitive to the various MSSM parameters, both in the mSUGRA and rSUGRA models, as may have been expected. However, these plots of the deviation of the observables from their respective SM values is more pronounced for the rSUGRA model than the mSUGRA model. This is essentially due to the additional parameter in the rSUGRA model, which is controlled by the pseudo-scalar Higgs mass. The value of the new Wilson coefficients corresponding to scalar and pseudo-scalar operators is directly proportional to $\tan^3\beta$ and m_A^2 and lower the value of m_A the higher the value of those Wilsons. Similarly, the higher the $\tan\beta$ the higher the new Wilsons. We effectively have that for the rSUGRA model a low m_A and high $\tan\beta$ region of MSSM parameter space also available which can generate large values of the new Wilsons. This was not the case in the mSUGRA model.

In figure 2 we have plotted of the integrated branching ratio in the mSUGRA. The integrated branching ratio can, at best, be 4 to 5 times the SM branching ratio for low values of $m_{1/2}$ and very high values of $\tan\beta$. However, for the rSUGRA model there can be a much higher enhancement to the branching ratio (as given in figure 3) when compared to the SM value for low m_A and high $\tan\beta$. As can be seen from the other graphs of the averaged observables the various observables show marked deviation from their respective SM values for a very wide region of the MSSM parameter space. This deviation cannot be explained solely on the basis of uncertainties in the form factors (which are at worst 15%). Thus such variations, if observed in future B-factories, could be very useful in testing the underlying operator structure of the theory and in fixing the numerical value of the Wilson coefficients.

There is also a further aspect to our results in relation to CP asymmetry. Consider the FB asymmetry of the conjugate process $b \rightarrow s\tau^+\tau^-$, where due to the smallness of the coupling of the b -quark with the u -quark the CKM-factor in all amplitudes involving the $b \rightarrow s$ transition, like the present one, will essentially be an overall factor. In the version of supersymmetry that we have considered and the parameter space thereof, there are no extra CP-violating phases. Thus in calculating decay rates the phase will be washed away and we have in effect a CP-invariant theory. The asymmetries of the process $b \rightarrow s\tau^-\tau^+$ and the conjugate process $\bar{b} \rightarrow \bar{s}\tau^+\tau^-$ are thus related. In fact the unpolarized FB asymmetry for the $b \rightarrow s$ transition will vanish in an untagged (CP even) sample [25]. Defining the forward and backward directions as referring to the τ^- and denoting the asymmetries of the conjugate process by \bar{A} , we get:

$$\bar{A}_{ij} = p_{ij} A_{ji} \quad (4.3)$$

with the parity factor p_{ij} equaling -1 for all i 's and j 's except for the combinations LN and NT . If we have an untagged sample containing an equal number of B 's and \bar{B} 's, then just as in the unpolarized sample, the asymmetries observed for the combinations LL , NN , TT as well as for $(LT + TL)$ will vanish. However, for the combinations $(LN + NL)$ and $(NT + TN)$, the asymmetries will add up. We thus have a situation in these two cases wherein a measurement of FB asymmetry for an untagged sample can lead to a meaningful non-null valued comparison between theory and experiment.

	f_+	f_0	f_T
$F(0)$	0.319	0.319	0.355
c_1	1.465	0.633	1.478
c_2	0.372	- 0.095	0.373
c_3	0.782	0.591	0.700

TABLE II: Form factors for $B \rightarrow K$ transition

APPENDIX A: FORM FACTORS

The form factors for the $B \rightarrow K$ transition are given in reference [2]:

$$\langle K(p_K) | \bar{s} \gamma_\mu b | B(p_B) \rangle = f_+ \left[(p_B + p_K)_\mu - \frac{m_B^2 - m_K^2}{q^2} q_\mu \right] + f_0 \frac{m_B^2 - m_K^2}{q^2} q_\mu \quad (\text{A1})$$

$$\langle K(p_K) | \bar{s} \sigma_{\mu\nu} q^\nu b | B(p_B) \rangle = i \frac{f_T}{m_B + m_K} \left[(p_B + p_K)_\mu q^2 - q_\mu (m_B^2 - m_K^2) \right] \quad (\text{A2})$$

where $q (= p_+ + p_-)$ is the sum of four momentas ℓ^- and ℓ^+ , i.e. the momentum transfer; f_+ , f_0 and f_T are the form factors. Multiplying equation (A1) by q^μ and by using the equations of motion we get:

$$\langle K(p_K) | \bar{s} b | B(p_B) \rangle = f_0 \frac{(m_B^2 - m_K^2)}{m_b} \quad (\text{A3})$$

where all the other matrix elements vanish.

For the form factors, f_+ , f_0 and f_T we will take the parameterization:

$$F(\hat{s}) = F(0) \exp(c_1 \hat{s} + c_2 \hat{s}^2 + c_3 \hat{s}^3) \quad (\text{A4})$$

where the values of the parameters are given in Table II

APPENDIX B: INPUT PARAMETERS

$$\begin{aligned} m_B &= 5.26 \text{GeV} , \quad m_b = 4.8 \text{GeV} , \quad m_c = 1.4 \text{GeV} \\ m_\mu &= 0.106 \text{GeV} , \quad m_\tau = 1.77 \text{GeV} , \quad m_w = 80.4 \text{GeV} , \quad m_z = 91.19 \text{GeV} \\ V_{tb} V_{ts}^* &= 0.0385 , \quad \alpha = \frac{1}{129} , \quad m_K = 0.49 \text{GeV} , \quad \Gamma_B = 4.22 \times 10^{-13} \text{GeV} \\ G_F &= 1.17 \times 10^{-5} \text{GeV}^{-2} \end{aligned}$$

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- [1] G. Buchalla, A. J. Buras and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996) [arXiv:hep-ph/9512380].
- [2] A. Ali, P. Ball, L. T. Handoko and G. Hiller, Phys. Rev. D **61**, 074024 (2000) [arXiv:hep-ph/9910221].
- [3] CLEO collaboration, T.E.Cowan *et al.* Phys. Rev. Lett. **84**, 5283 (2000) [arXiv:hep-ex/9908022] ; ALEPH Collaboration, R.Barate *et al.* Phys. Lett. B **429**, 169 (1998) .
- [4] J. L. Lopez, D. V. Nanopoulos, X. Wang and A. Zichichi, Phys. Rev. D **51**, 147 (1995) [arXiv:hep-ph/9406427] ; R. Barbieri and G. F. Giudice, Phys. Lett. B **309**, 86 (1993) [arXiv:hep-ph/9303270] ; T. Goto and Y. Okada, Prog. Theor. Phys. **94**, 407 (1995) [arXiv:hep-ph/9412225].

- [5] W. Bensalem, D. London, N. Sinha and R. Sinha, Phys. Rev. **D 67**, 034007 (2003) [arXiv:hep-ph/0209228]; Naveen Gaur, arXiv : hep-ph/0305242 ; A. S. Cornell and N. Gaur, JHEP **0309**, 030 (2003) [arXiv:hep-ph/0308132].
- [6] S. Rai Choudhury, Naveen Gaur, A. S. Cornell and G. C. Joshi, Phys. Rev. **D 68**, 054016 (2003) [arXiv:hep-ph/0304084].
- [7] F. Krüger and L. M. Sehgal, Phys. Lett. B **380**, 199 (1996), [arXiv:hep-ph/9603237] ; J. L. Hewett, Phys. Rev. D **53**, 4964 (1996), [arXiv:hep-ph/9506289]; S. Rai Choudhury, A. Gupta and N. Gaur, Phys. Rev. D **60**, 115004 (1999)[arXiv:hep-ph/9902355] ; S. Fukae, C. S. Kim and T. Yoshikawa, Phys. Rev. D **61**, 074015 (2000)[arXiv:hep-ph/9908229] ;
- [8] C. Q. Geng and C. P. Kao, Phys. Rev. D **57**, 4479 (1998). T. M. Aliev, M. K. Cakmak and M. Savci, Nucl. Phys. B **607**, 305 (2001) [arXiv:hep-ph/0009133] ; T. M. Aliev and M. Savci, Phys. Lett. B **481**, 275 (2000) [arXiv:hep-ph/0003188]. T. M. Aliev, C. S. Kim and Y. G. Kim, Phys. Rev. D **62**, 014026 (2000) [arXiv:hep-ph/9910501].
- [9] T. M. Aliev, M. K. Cakmak, A. Ozpineci and M. Savci, Phys. Rev. D **64**, 055007 (2001) [arXiv:hep-ph/0103039] ; T. M. Aliev, M. Savci, A. Ozpineci and H. Koru, J. Phys. G **24**, 49 (1998) [arXiv:hep-ph/9705222]. T. M. Aliev, M. K. Cakmak, A. Ozpineci and M. Savci, Phys. Rev. D **64**, 055007 (2001) [arXiv:hep-ph/0103039] ;
- [10] T. M. Aliev, A. Ozpineci, M. Savci and C. Yuce, Phys. Rev. D **66**, 115006 (2002) [arXiv:hep-ph/0208128] ; T. M. Aliev, A. Ozpineci and M. Savci, Phys. Lett. B **511**, 49 (2001) [arXiv:hep-ph/0103261] ; T. M. Aliev, A. Ozpineci and M. Savci, Nucl. Phys. B **585**, 275 (2000) [arXiv:hep-ph/0002061] ; T. M. Aliev, D. A. Demir and M. Savci, Phys. Rev. D **62**, 074016 (2000) [arXiv:hep-ph/9912525]; G. Erkol and G. Turan, Nucl. Phys. B **635**, 286 (2002) [arXiv:hep-ph/0204219].
- [11] Q. S. Yan, C. S. Huang, W. Liao and S. H. Zhu, Phys. Rev. D **62**, 094023 (2000) [arXiv:hep-ph/0004262].
- [12] C. Bobeth, T. Ewerth, F. Kruger and J. Urban, Phys. Rev. D **64**, 074014 (2001) [arXiv:hep-ph/0104284]. D. A. Demir, K. A. Oliev and M. B. Voloshin Phys. Rev. D **66**, 034015 (2002) [arXive : hep-ph/0204119] ;
- [13] E. O. Iltan , Int. J. Mod. Phys. **A 14**, 4365 (1999), [arXive : hep-ph/9807256] ; T. M. Aliev and M. Savci , Phys. Rev. **D 60**, 014005 (1999), [arXive : hep-ph/9812272] ; S. R. Choudhury and N. Gaur, Phys. Rev. D **66**, 094015 (2002) [arXiv:hep-ph/0206128] ; G. Erkol and G. Turan, JHEP **0202**, 015 (2002) [arXiv:hep-ph/0201055].
- [14] S. R. Choudhury and N. Gaur, Phys. Lett. B **451**, 86 (1999), [arXiv:hep-ph/9810307] ; P. H. Chankowski and L. Slawianowska, Phys. Rev. D **63**, 054012 (2001) [arXiv:hep-ph/0008046]; A. J. Buras, P. H. Chankowski, J. Rosiek and L. Slawianowska, Nucl. Phys. B **659**, 3 (2003) [arXiv:hep-ph/0210145]; A. J. Buras, P. H. Chankowski, J. Rosiek and L. Slawianowska, Phys. Lett. B **546**, 96 (2002) [arXiv:hep-ph/0207241]; J. K. Mizukoshi, X. Tata and Y. Wang, Phys. Rev. D **66**, 115003 (2002) [arXiv:hep-ph/0208078]; K. S. Babu and C. F. Kolda, Phys. Rev. Lett. **84**, 228 (2000) [arXiv:hep-ph/9909476] ; T. Ibrahim and P. Nath, Phys. Rev. D **67**, 016005 (2003) [arXiv:hep-ph/0208142] ; C. S. Huang and W. Liao, Phys. Lett. B **538**, 301 (2002) [arXiv:hep-ph/0201121]; S. Baek, P. Ko and W. Y. Song, Phys. Rev. Lett. **89**, 271801 (2002) [arXiv:hep-ph/0205259] ; S. Baek, P. Ko and W. Y. Song, JHEP **0303**, 054 (2003) [arXiv:hep-ph/0208112]; A. Dedes and A. Pilaftsis, Phys. Rev. D **67**, 015012 (2003) [arXiv:hep-ph/0209306] ; C. S. Huang, W. Liao, Q. S. Yan and S. H. Zhu, Phys. Rev. D **63**, 114021 (2001) [Erratum-ibid. D **64**, 059902 (2001)] [arXiv:hep-ph/0006250]; C. S. Huang and X. H. Wu, Nucl. Phys. B **657**, 304 (2003) [arXiv:hep-ph/0212220].
- [15] W. Skiba and J. Kalinowski, Nucl. Phys. B **404**, 3 (1993) ; H. E. Logan and U. Nierste, Nucl. Phys. B **586**, 39 (2000), [arXiv:hep-ph/0004139] ;
- [16] S. Rai Choudhury, N. Gaur and N. Mahajan, Phys. Rev. D **66**, 054003 (2002) [arXiv:hep-ph/0203041] ; S. R. Choudhury and N. Gaur, arXiv:hep-ph/0205076 ; E. O. Iltan and G. Turan, Phys. Rev. D **61**, 034010 (2000), [arXiv:hep-ph/9906502] ; G. Erkol and G. Turan Acta. Phys. Pol. **B 33**, 1285, (2002) [arXiv:hep-ph/0112115] ; G. Erkol and G. Turan, Phys. Rev. D **65**, 094029 (2002), [arXiv:hep-ph/0110017] ; T. M. Aliev, A. Ozpineci, M. Savci, Phys. Lett. B **520**, 69 (2001), [arXiv:hep-ph/0105279] .
- [17] Z. Xiong and J. M. Yang, Nucl. Phys. B **628**, 193 (2002)[arXiv:hep-ph/0105260] ; C. Bobeth, A. J. Buras, F. Kruger and J. Urban, Nucl. Phys. B **630**, 87 (2002) [arXiv:hep-ph/0112305] ; C. Huang, W. Liao and Q. Yan, Phys. Rev. D **59**, 011701 (1999), [arXiv:hep-ph/9803460].
- [18] B. Grinstein, M. J. Savage and M. B. Wise, Nucl. Phys. B **319**, 271 (1989) ; A. J. Buras and M. Münz, Phys. Rev. D **52**, 186 (1995) [arXiv:hep-ph/9501281].
- [19] P. L. Cho, M. Misiak and D. Wyler, Phys. Rev. D **54**, 3329 (1996) [arXiv:hep-ph/9601360]; J. L. Hewett and J. D. Wells, Phys. Rev. D **55**, 5549 (1997) [arXiv:hep-ph/9610323].
- [20] A. Ali, T. Mannel and T. Morozumi, Phys. Lett. B **273**, 505 (1991); C. S. Lim, T. Morozumi and A. I. Sanda, Phys. Lett. B **218**, 343 (1989); N. G. Deshpande, J. Trampetic and K. Panose, Phys. Rev. D **39**, 1461 (1989); P. J. O'Donnell and H. K. Tung, Phys. Rev. D **43**, 2067 (1991) .
- [21] A. Ali, T. Mannel and T. Morozumi, Phys. Lett. B **273**, 505 (1991).
- [22] H. P. Nilles, Phys. Rept. **110**, 1 (1984) ; H. E. Haber and G. L. Kane, Phys. Rept. **117**, 75 (1985).
- [23] T. Goto, Y. Okada, Y. Shimizu and M. Tanaka, Phys. Rev. D **55**, 4273 (1997) [arXiv:hep-ph/9609512] ; T. Goto, Y. Okada and Y. Shimizu, Phys. Rev. D **58**, 094006 (1998) [arXiv:hep-ph/9804294] ; J. R. Ellis, K. A. Olive and Y. Santoso, arXiv : hep-ph/0204192.
- [24] B Physics at Tevatron : Run II & Beyond , K. Anikeev *et al.* arXiv : hep-ph/0201071 .
- [25] S. R. Choudhury, Phys. Rev. D **56**, 6028 (1997) [arXiv:hep-ph/9706313]. G. Buchalla, G. Hiller and G. Isidori, Phys. Rev. D **63**, 014015 (2001) [arXiv:hep-ph/0006136].
- [26] M. Beneke and M. Neubert, arXiv:hep-ph/0308039.
- [27] A. H. Chamseddine, R. Arnowitt and P. Nath, Phys. Rev. Lett. **49**, 970 (1982) ; L. J. Hall, J. Lykken and S. Weinberg, Phys. Rev. D **27**, 2359 (1983) ; R. Barbieri, S. Ferrara and C. A. Savoy, Phys. Lett. B **119**, 343 (1982).
- [28] L. Randall and R. Sundrum, Nucl. Phys. B **557**, 79 (1999) [arXiv:hep-th/9810155] ; G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP **9812**, 027 (1998) [arXiv:hep-ph/9810442] ; J. A. Bagger, T. Moroi and E. Poppitz, JHEP

- 0004**, 009 (2000) [arXiv:hep-th/9911029] ; K. Huitu, J. Laamanen and P. N. Pandita, Phys. Rev. D **65**, 115003 (2002) [arXiv:hep-ph/0203186].
- [29] G. F. Giudice and R. Rattazzi, Phys. Rept. **322**, 419 (1999) [arXiv:hep-ph/9801271].
- [30] M. Battaglia *et al.*, Eur. Phys. J. C **22**, 535 (2001) [arXiv:hep-ph/0106204] ; J. R. Ellis, K. A. Olive and Y. Santoso, Phys. Lett. B **539**, 107 (2002) [arXiv:hep-ph/0204192].
- [31] K. Hagiwara *et al.* [Particle Data Group Collaboration], Phys. Rev. D **66**, 010001 (2002).

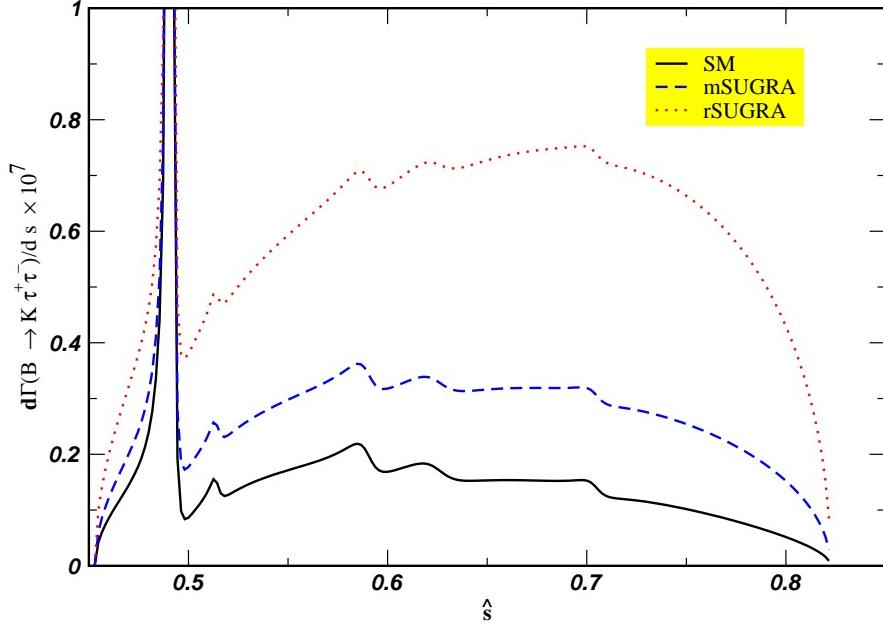


FIG. 1: The differential decay rate $d\Gamma/d\hat{s}$ against \hat{s} . The parameters for the mSUGRA model are: $m_0 = 400\text{GeV}$, $m_{1/2} = 500\text{GeV}$, $\tan\beta = 40$ and $A = 0$. The additional parameter for the rSUGRA models is taken to be $m_A = 345\text{GeV}$.

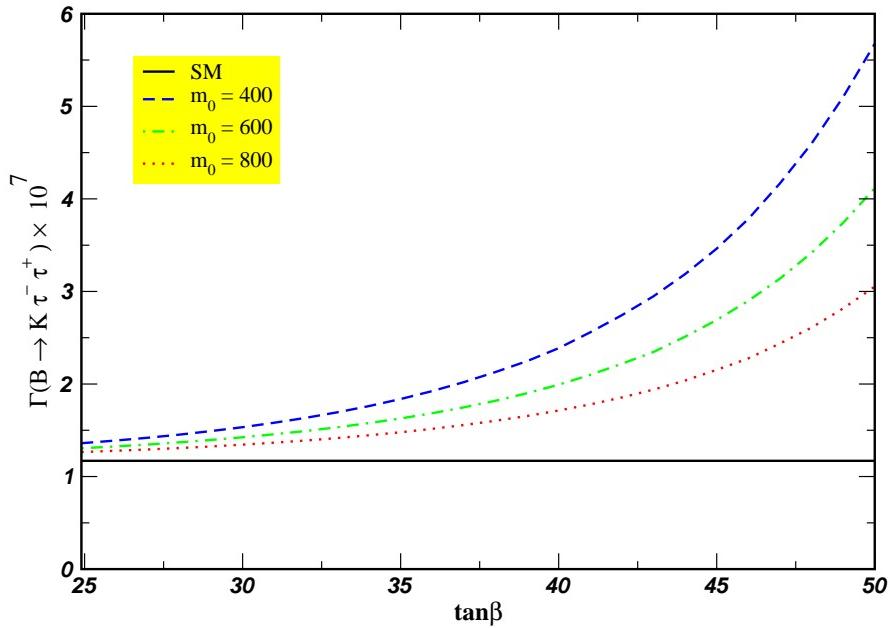


FIG. 2: $\text{Br}(B \rightarrow K\tau^-\tau^+)$ variation with $\tan\beta$ in mSUGRA for various values of m_0 . Other model parameters are: $m_{1/2} = 500\text{GeV}$, $A = 0$.

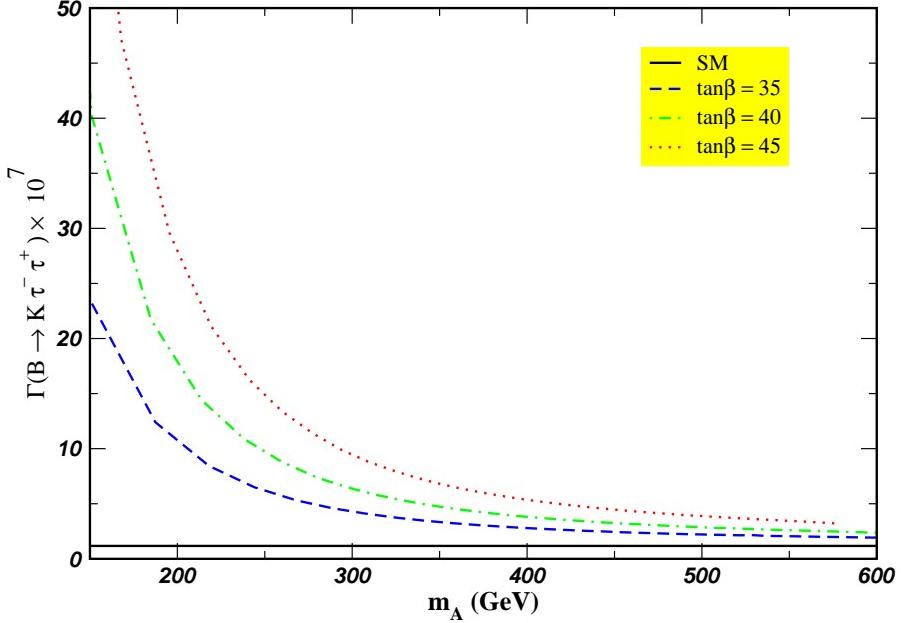


FIG. 3: $\text{Br}(B \rightarrow K\tau^-\tau^+)$ variation with m_A in rSUGRA for various $\tan\beta$ values. Other model parameters: $m_0 = 500\text{GeV}$, $m_{1/2} = 500\text{ GeV}$, $A = 0$.

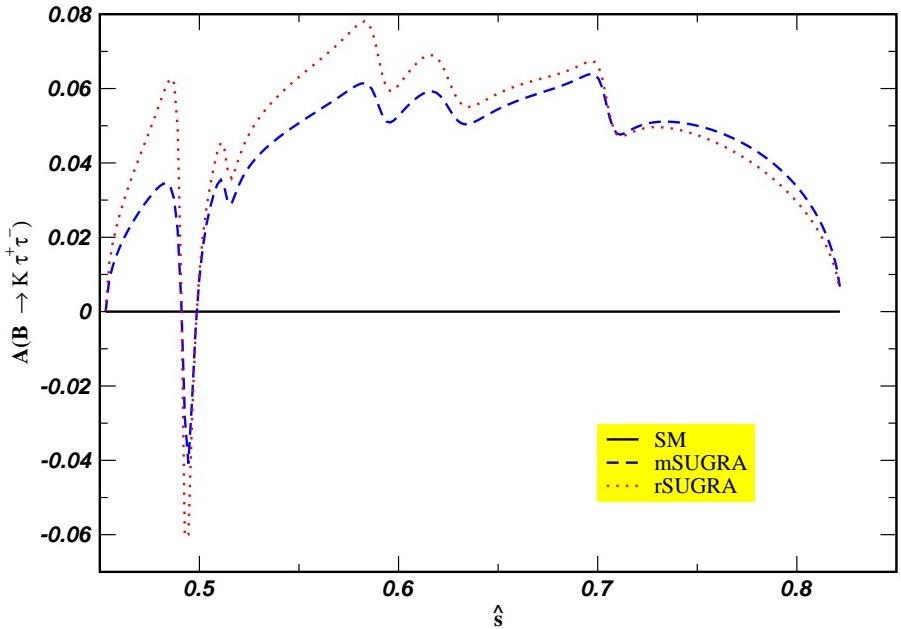


FIG. 4: Unpolarized FB asymmetry with dilepton invariant mass \hat{s} . Other parameters of mSUGRA and rSUGRA are given in figure 1.

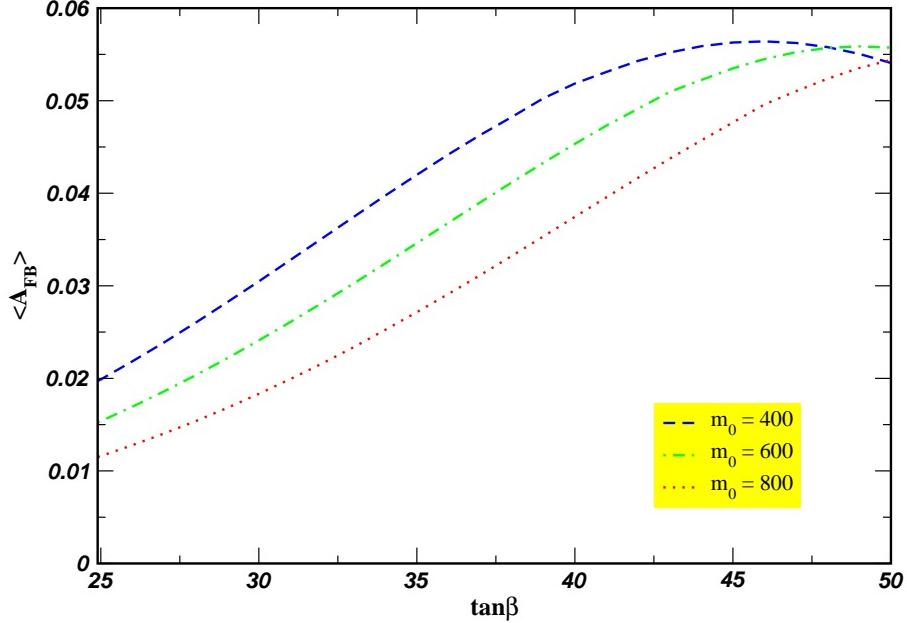


FIG. 5: Integrated FB asymmetry variation with $\tan\beta$ in mSUGRA model. Other parameters are same as given in figure 2.

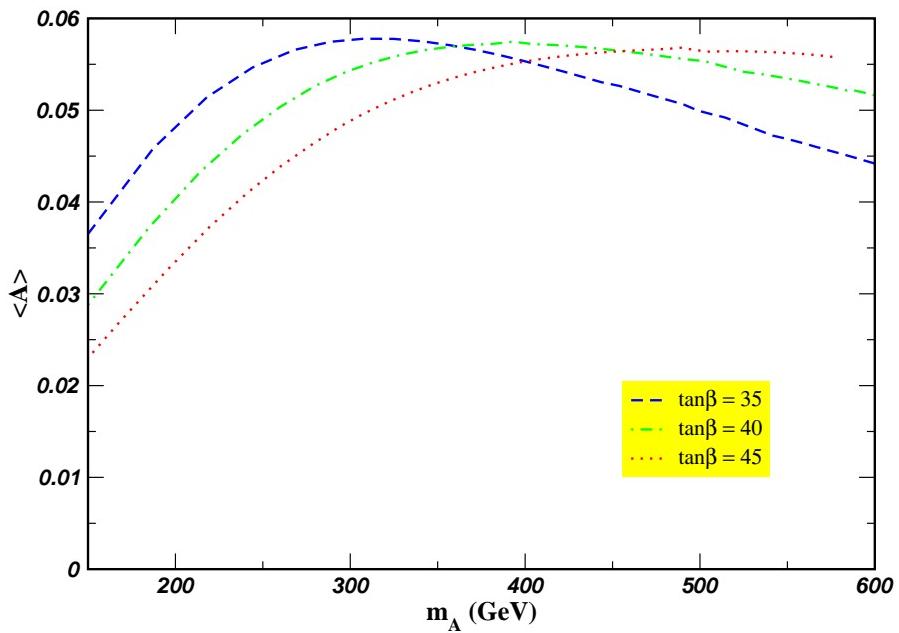


FIG. 6: Integrated FB asymmetry variation with m_A in rSUGRA model, with other model parameters as given in figure 3.

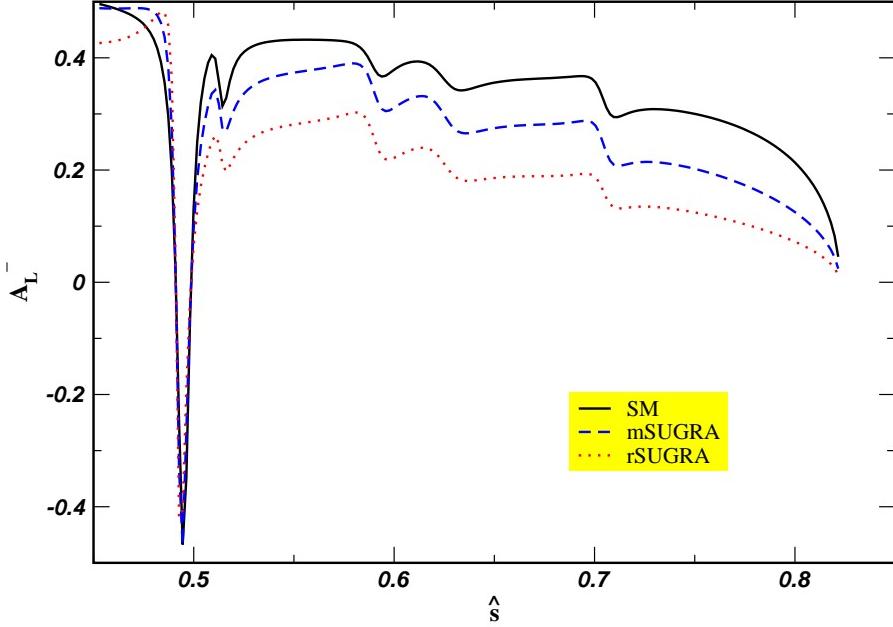


FIG. 7: A_L^- variation with \hat{s} . Other parameters of mSUGRA and rSUGRA model are the same as given in figure 1.

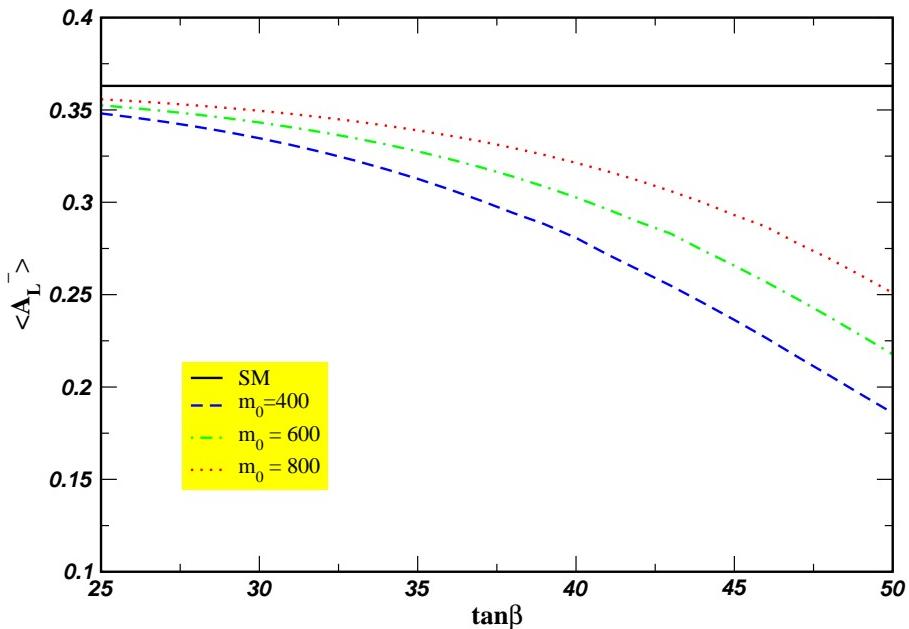


FIG. 8: $\langle A_L^- \rangle$ variation with $\tan\beta$ in mSUGRA model for various values of m_0 . Other model parameters are as given in figure 2.

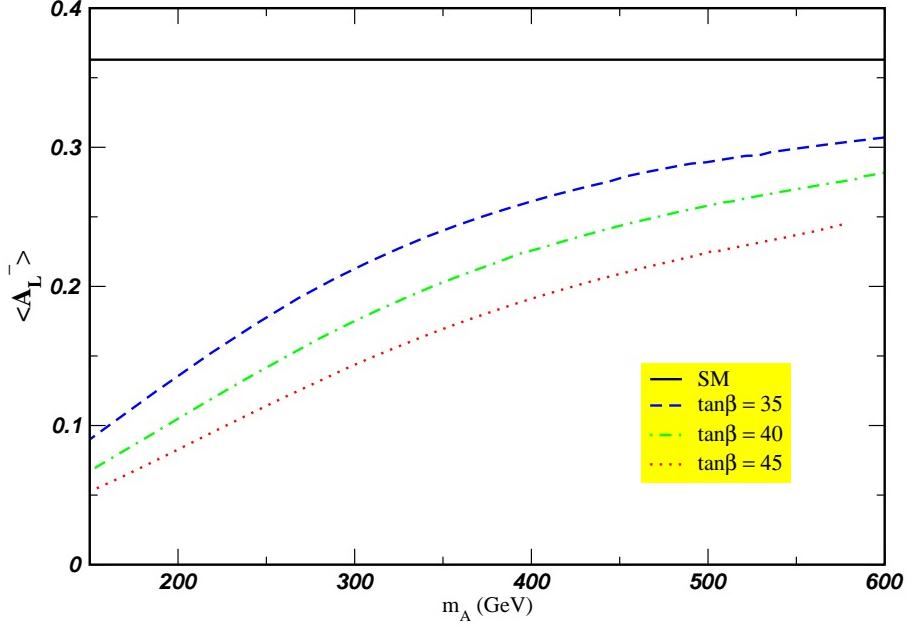


FIG. 9: $\langle A_L^- \rangle$ variation with m_A in rSUGRA model for various values of $\tan\beta$. Other model parameters are as given in figure 3.

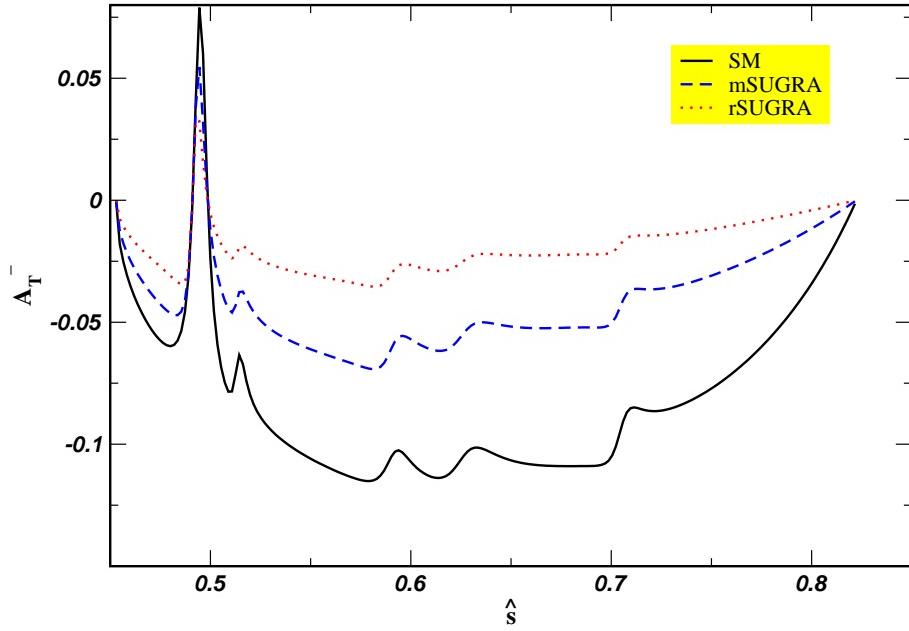


FIG. 10: A_T^- with \hat{s} . Other parameters are the same as given in figure1.

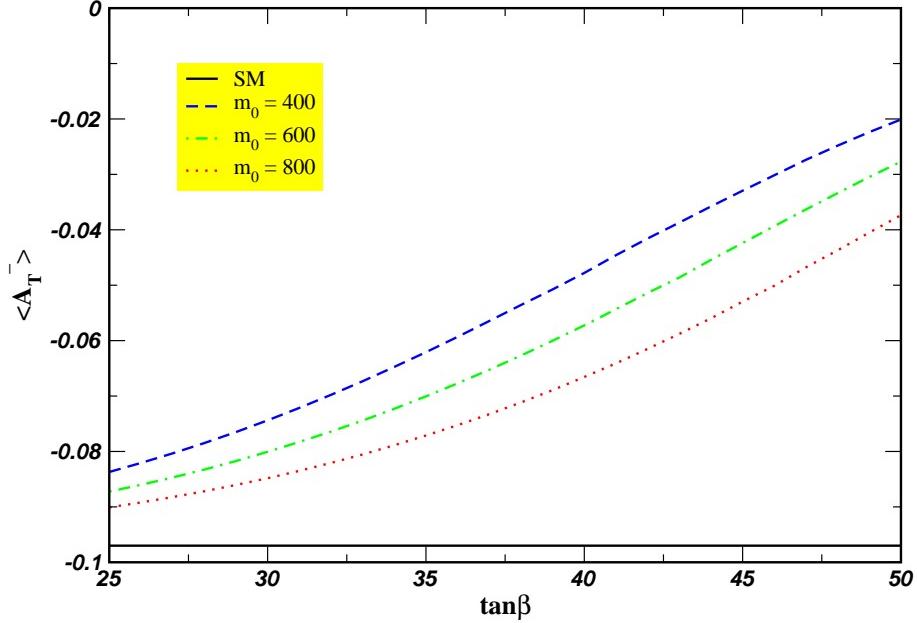


FIG. 11: $\langle A_T^- \rangle$ with $\tan\beta$ in mSUGRA, with other parameters same as in figure 9.

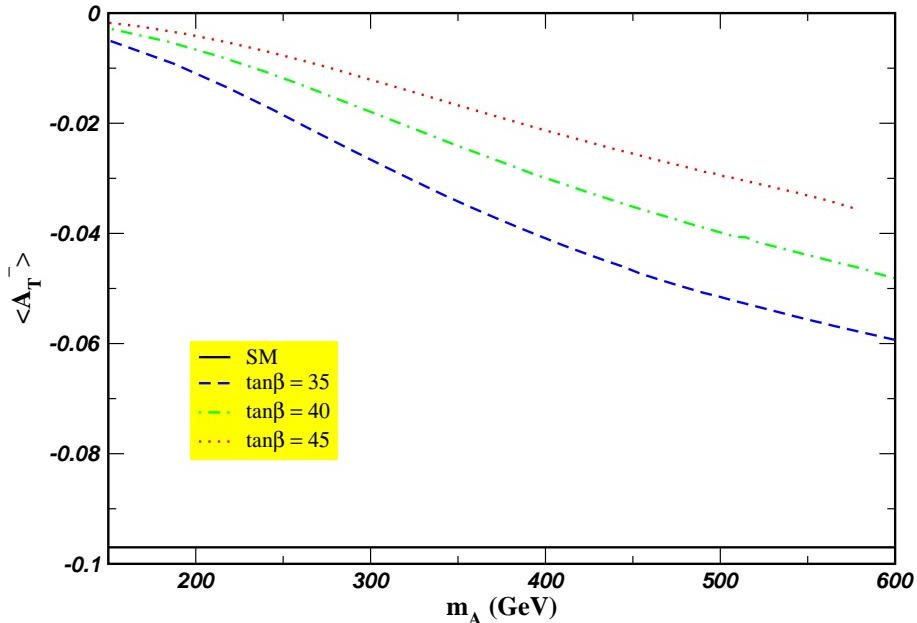


FIG. 12: $\langle A_T^- \rangle$ with m_A in rSUGRA, with other parameters same as in figure 10.

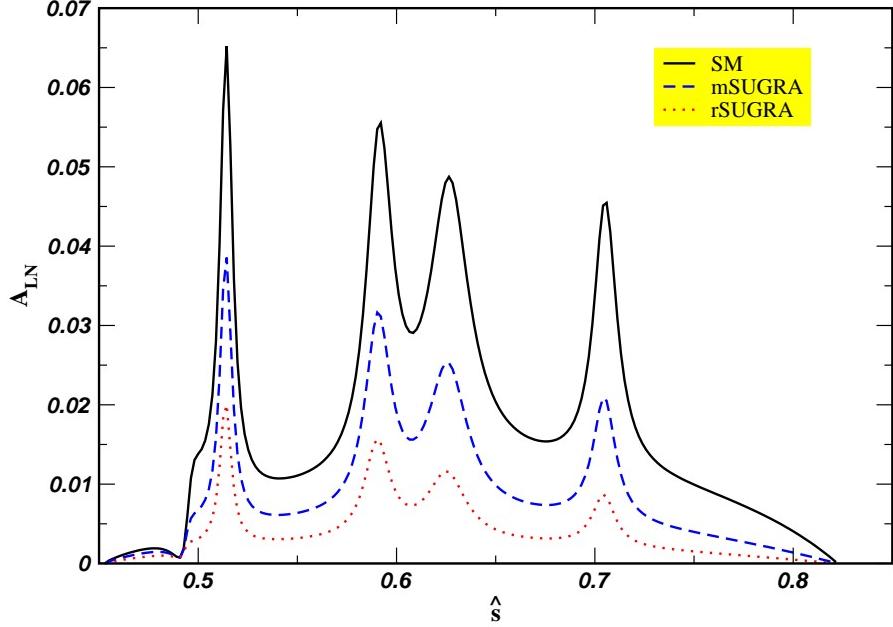


FIG. 13: A_{LN} with \hat{s} . Other parameters are the same as given in figure 1.

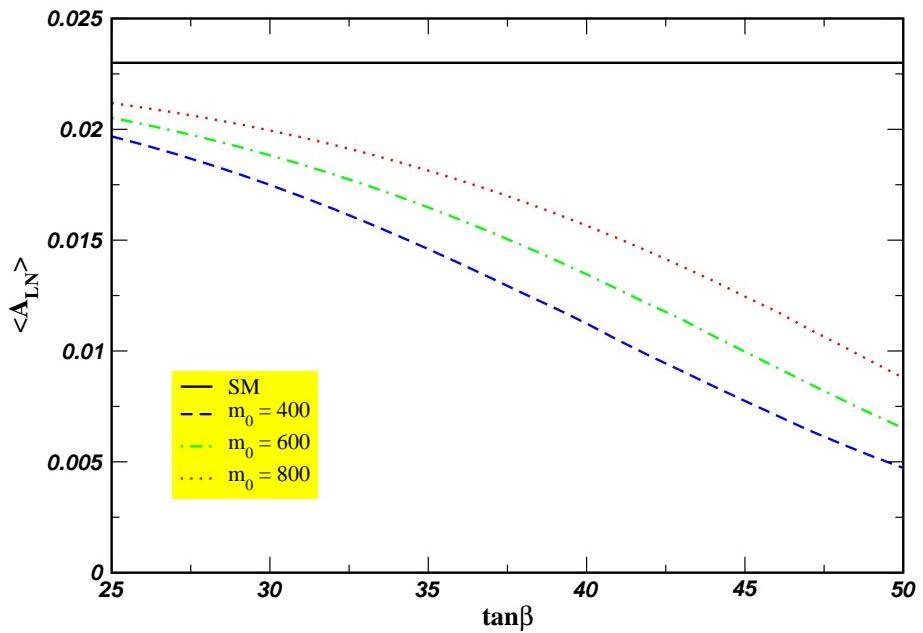


FIG. 14: $\langle A_{LN} \rangle$ with $\tan\beta$ in mSUGRA, with other parameters Same as in figure 9.

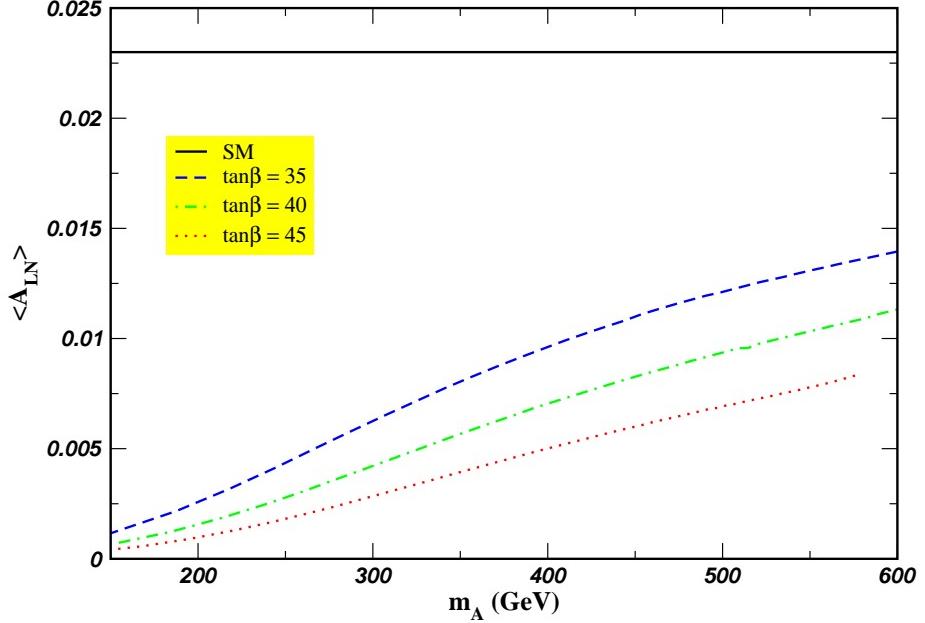


FIG. 15: $\langle A_{LN} \rangle$ with m_A in rSUGRA model, with other parameters same as in figure 10.

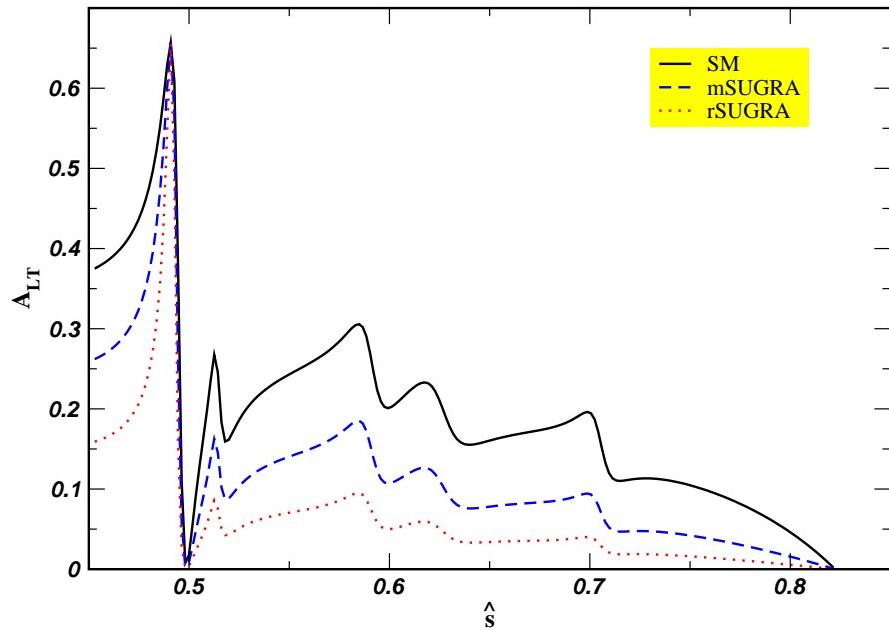


FIG. 16: A_{LT} with \hat{s} . Other parameters are the same as given in figure 1.

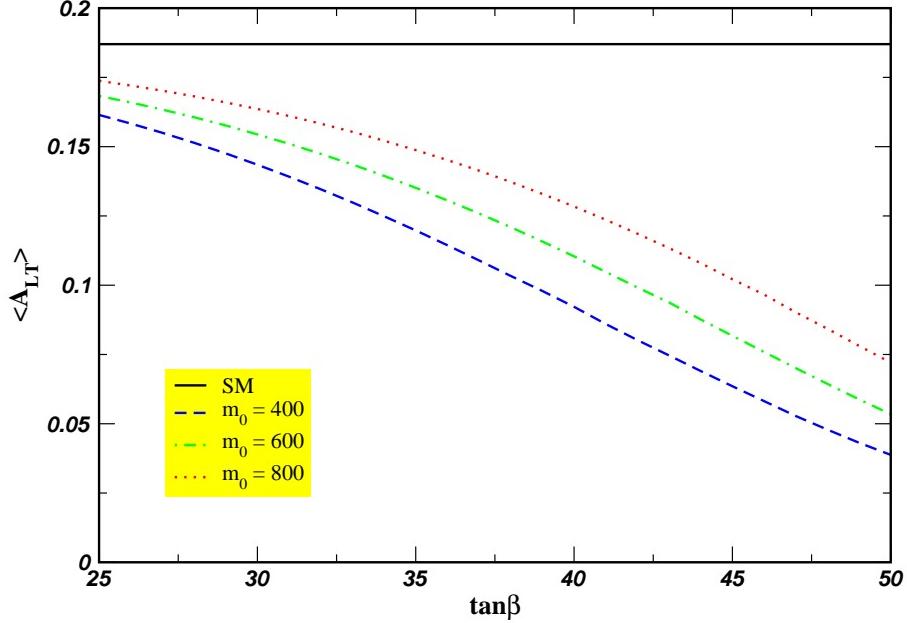


FIG. 17: $\langle A_{LT} \rangle$ with $\tan\beta$ in mSUGRA, with the other parameters being the same as in figure 9.

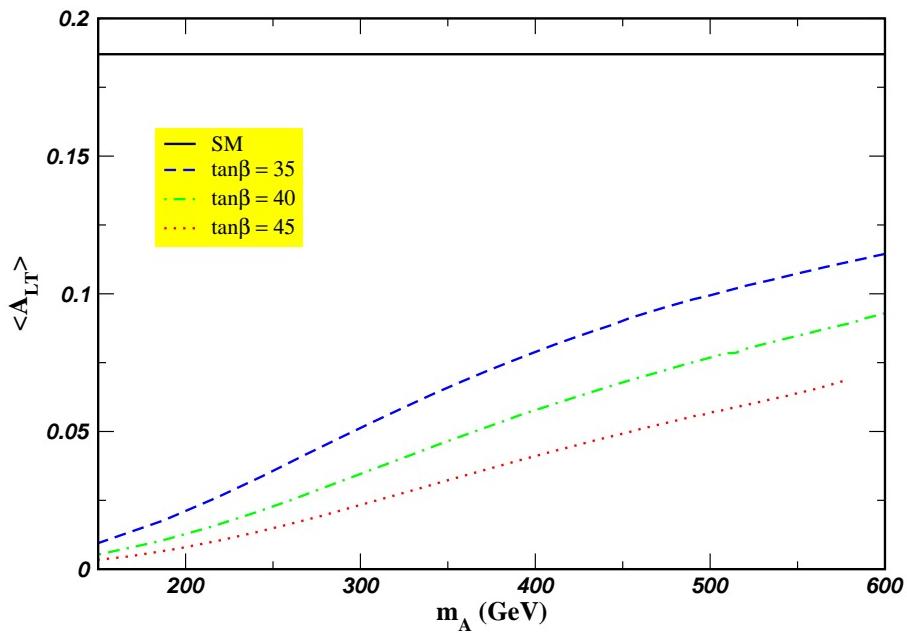


FIG. 18: $\langle A_{LT} \rangle$ with m_A , with the other parameters the same as in figure 10.

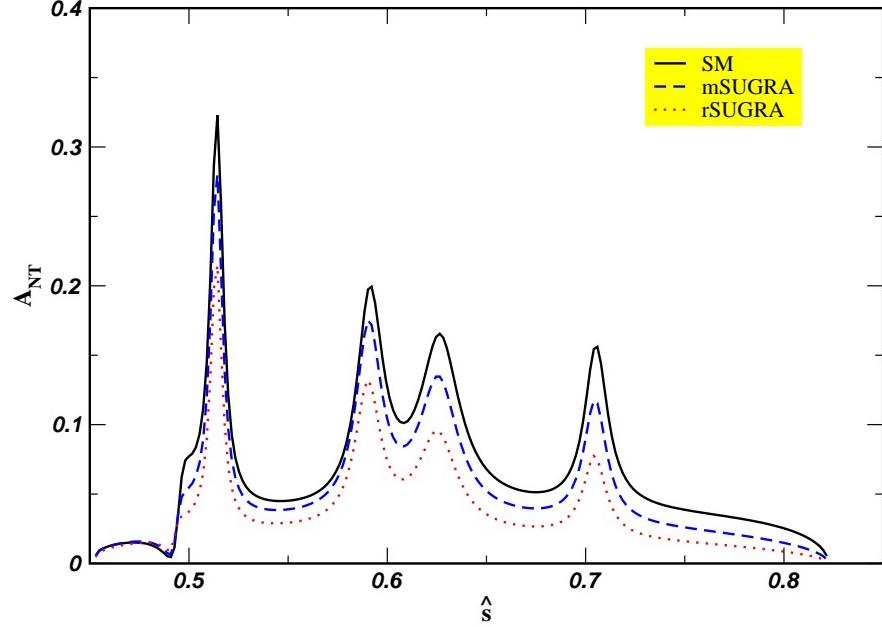


FIG. 19: A_{NT} with \hat{s} . Other parameters are the same as given in figure 1.

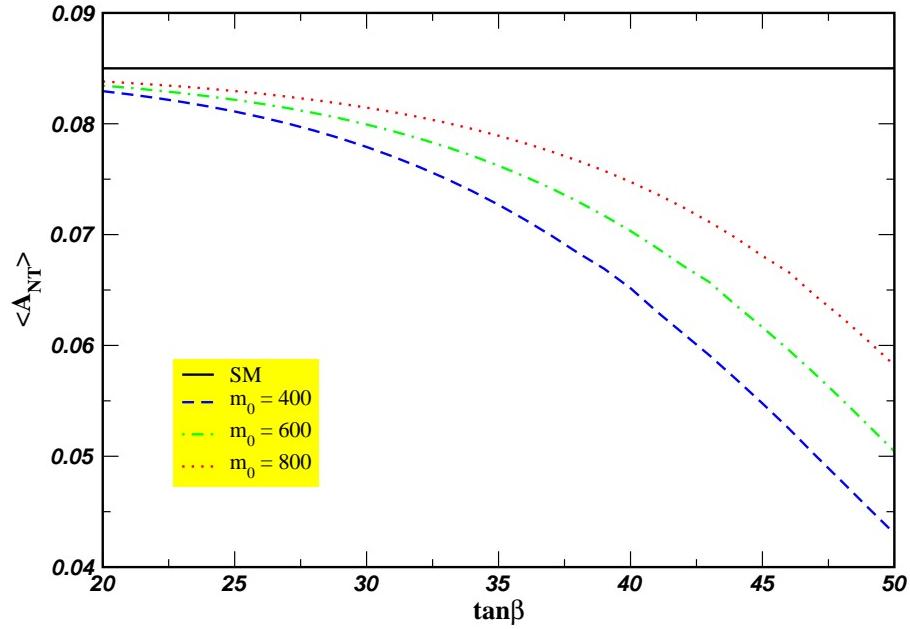


FIG. 20: $\langle A_{NT} \rangle$ with $\tan\beta$ in mSUGRA, with other parameters same as in figure 9.

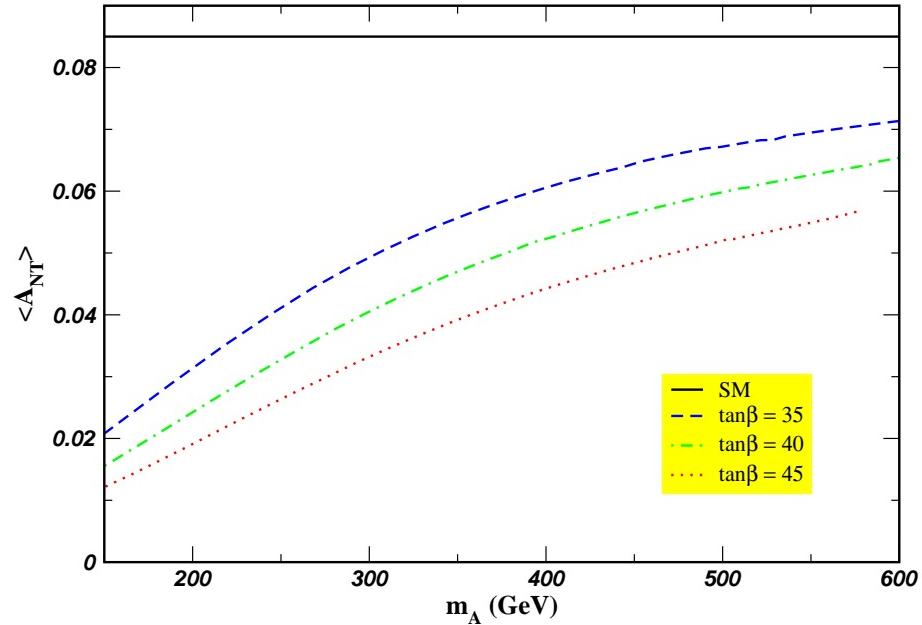


FIG. 21: $\langle A_{NT} \rangle$ with m_A in rSUGRA model, with the other parameters the same as in figure 10.